

Accessing transversity GPDs in neutrino-production of a charmed meson*

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We demonstrate that the transversity chiral odd generalized parton distributions can be accessed through the azimuthal dependence of the $\nu_l N \rightarrow l^- D^+ N'$ or $\bar{\nu}_l N \rightarrow l^+ D^- N'$ differential cross sections, in the framework of the collinear QCD approach, where GPDs factorize from perturbatively calculable coefficient functions calculated up to order m_c/Q .

PACS numbers: 14.20.Dh, 12.38.Bx, 13.15.+g

1. Introduction

Thanks to quark mass effects, flavor changing electroweak interactions open new ways to access chiral-odd quantities like the transversity generalized parton distributions (GPDs) in a nucleon [1]. The transversity distributions which encode the transverse spin structure of the nucleon have proven to be among the most difficult hadronic quantities to access. This is due to the chiral odd character of the quark operators which enter their definition; this feature enforces the decoupling of these distributions from most measurable hard amplitudes. After the pioneering works [2], much effort [3] has been devoted to the exploration of many channels but experimental difficulties have challenged some of the most promising ones.

Generalized parton distributions give access to the internal structure of hadrons in a much more detailed way than parton distributions measured in inclusive processes, since they allow a 3-dimensional analysis [4]. The study of exclusive reactions mediated by a highly virtual photon in the

* Presented at 16th conference on Elastic and diffractive scattering, EDS Blois 2015, Borgo, Corsica, France, June 29th-July 4th, 2015

generalized Bjorken regime benefits from the factorization properties of the leading twist QCD amplitudes [5, 6, 7] of reactions such as deeply virtual Compton scattering.

Neutrino production offers another way to access (generalized) parton distributions [8]. Although neutrino induced cross sections are orders of magnitudes smaller than those for electroproduction, the advent of new generations of neutrino experiments will open new possibilities. We want here to stress that they can help to access the elusive [9] chiral-odd generalized parton distributions.

2. The processes

We consider [1] the exclusive reactions

$$\begin{aligned}\nu_l(k)N(p_1) &\rightarrow l^-(k')D^+(p_D)N'(p_2), \\ \bar{\nu}_l(k)N(p_1) &\rightarrow l^+(k')D^-(p_D)N'(p_2),\end{aligned}\tag{1}$$

in the kinematical domain where collinear factorization leads to a description of the scattering amplitude in terms of nucleon GPDs and the D -meson distribution amplitude, with the hard subprocess ($q = k' - k; Q^2 = -q^2$):

$$W^+(q)d \rightarrow D^+d' \quad \text{or} \quad W^-(q)u \rightarrow D^-u',\tag{2}$$

described for the neutrino case by the handbag Feynman diagrams of Fig. 1. We use the standard notations of deep exclusive leptonproduction, namely $P = (p_1 + p_2)/2$, $\Delta = p_2 - p_1$, $t = \Delta^2$, $x_B = Q^2/2p_1 \cdot q$, $y = p_1 \cdot q/p_1 \cdot k$ and $\epsilon \simeq 2(1-y)/[1+(1-y)^2]$. p and n are light-cone vectors ($v \cdot n = v^+$, $v \cdot p = p^-$ for any vector v) and $\xi = -\Delta \cdot n/2P \cdot n$ is the skewness variable.

2.1. Transversity GPDs.

The four twist 2 quark transversity GPDs have been defined [10] as :

$$\begin{aligned}&\frac{1}{2} \int \frac{dz^-}{2\pi} e^{ixP^+z^-} \langle p_2, \lambda' | \bar{\psi}(-\frac{1}{2}z) i\sigma^{+i} \psi(\frac{1}{2}z) | p_1, \lambda \rangle \Big|_{z^+ = \mathbf{z}_T = 0} \\ &= \frac{1}{2P^+} \bar{u}(p_2, \lambda') \left[H_T^q i\sigma^{+i} + \tilde{H}_T^q \frac{P^+ \Delta^i - \Delta^+ P^i}{m_N^2} \right. \\ &\quad \left. + E_T^q \frac{\gamma^+ \Delta^i - \Delta^+ \gamma^i}{2m_N} + \tilde{E}_T^q \frac{\gamma^+ P^i - P^+ \gamma^i}{m_N} \right] u(p_1, \lambda).\end{aligned}\tag{3}$$

and their experimental access is much discussed in [11]. The leading GPD $H_T(x, \xi, t)$ is equal to the transversity PDF in the $\xi = 0, t = 0$ limit. Models have been proposed [11, 12] and some lattice calculation results exist [13].

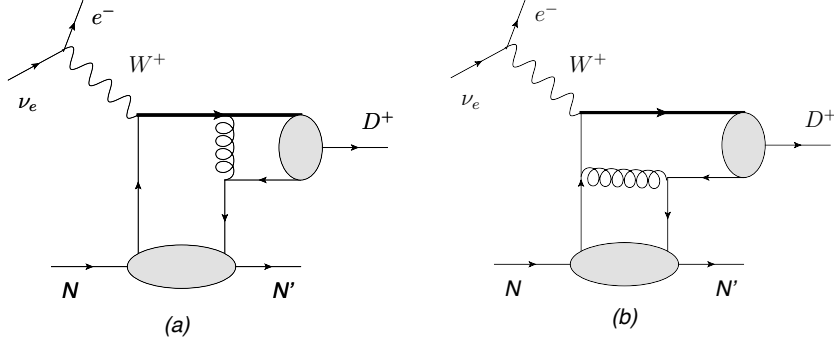


Fig. 1. Feynman diagrams for the factorized amplitude for the $\nu_e N \rightarrow e^- D^+ N'$ process; the thick line represents the heavy quark. In the Feynman gauge, diagram (a) involves convolution with both the transversity GPDs and the chiral even ones, whereas diagram (b) involves only chiral even GPDs.

2.2. The transverse amplitude up to $O(m_c/Q^2)$.

It has been demonstrated [14] that hard-scattering factorization of meson leptonproduction is valid at leading twist with the inclusion of heavy quark masses in the hard amplitude. In our case this means including the part $\frac{m_c}{k_c^2 - m_c^2}$ in the off-shell heavy quark propagator (see the Feynman graph on Fig. 1a) present in the leading twist coefficient function. We keep the term m_c^2 in the denominator since it will help us to understand precisely how to perform the integration over the longitudinal momentum fraction x around the point $x = \xi$. The longitudinal leading twist amplitude with a longitudinally polarized W boson T_L has been computed in [16]. Adding mass terms to the heavy quark propagator has no effect on the calculation of this amplitude, but leads to a non-zero transverse amplitude when a chiral-odd transversity GPD is involved.

In the Feynman gauge, the non-vanishing m_c -dependent part of the Dirac trace in the hard scattering part depicted in Fig. 1a reads:

$$\begin{aligned} \text{Tr} \left[\sigma^{pi} \gamma^\nu \hat{p}_D \gamma^5 \gamma_\nu \frac{m_c}{k_c^2 - m_c^2 + i\epsilon} (1 - \gamma^5) \hat{\epsilon} \frac{1}{k_g^2 + i\epsilon} \right] \\ = \frac{2Q^2}{\xi} \epsilon_\mu [i\epsilon^{\mu pin} - g_\perp^{\mu i}] \frac{m_c}{k_c^2 - m_c^2 + i\epsilon} \frac{1}{k_g^2 + i\epsilon}, \end{aligned} \quad (4)$$

where k_c (k_g) is the heavy quark (gluon) momentum and ϵ the polarization vector of the W -boson (we denote $\hat{p} = p_\mu \gamma^\mu$ for any vector p ; the transverse index i will be contracted with the analogous index in transversity GPDs). The fermionic trace vanishes for the diagram shown on Fig. 1b.

The transverse amplitude is then written as ($\tau = 1 - i2$):

$$T_T = \frac{iC\xi m_c}{\sqrt{2}Q^2} \bar{N}(p_2) \left[\mathcal{H}_T^\phi i\sigma^{n\tau} + \tilde{\mathcal{H}}_T^\phi \frac{\Delta^\tau}{m_N^2} + \mathcal{E}_T^\phi \frac{\hat{n}\Delta^\tau + 2\xi\gamma^\tau}{2m_N} - \tilde{\mathcal{E}}_T^\phi \frac{\gamma^\tau}{m_N} \right] N(p_1),$$

in terms of transverse form factors that we define as :

$$\mathcal{F}_T^\phi = f_D \int \frac{\phi(z)dz}{\bar{z}} \int \frac{F_T^d(x, \xi, t)dx}{(x - \xi + i\epsilon)(x - \xi + \alpha\bar{z} + i\epsilon)}, \quad (5)$$

where F_T^d is any d-quark transversity GPD, $\alpha = \frac{2\xi m_c^2}{Q^2 + m_c^2}$ and $\tilde{\mathcal{E}}_T^\phi = \xi\mathcal{E}_T^\phi - \tilde{\mathcal{E}}_T^\phi$. Keeping α in the denominator in Eq. (5) regulates the treatment of the double pole in the integration over x .

3. The angular dependence

The dependence of a leptonproduction cross section on azimuthal angles is widely used way to analyze the scattering mechanism, as for deeply virtual Compton scattering [15]. In the neutrino case, it reads:

$$\begin{aligned} \frac{d^4\sigma(\nu N \rightarrow l^- N' D)}{dx_B dQ^2 dt d\varphi} &= \tilde{\Gamma} \left\{ \frac{1 + \sqrt{1 - \varepsilon^2}}{2} \sigma_{--} + \varepsilon \sigma_{00} \right. \\ &\left. + \sqrt{\varepsilon}(\sqrt{1 + \varepsilon} + \sqrt{1 - \varepsilon})(\cos \varphi \operatorname{Re} \sigma_{-0} + \sin \varphi \operatorname{Im} \sigma_{-0}) \right\}, \end{aligned} \quad (6)$$

where the “cross-sections” $\sigma_{lm} = \epsilon_l^{*\mu} W_{\mu\nu} \epsilon_m^\nu$ are product of amplitudes for the process $W(\epsilon_l)N \rightarrow DN'$, averaged (summed) over the initial (final) hadron polarizations. The azimuthal angle φ is defined in the initial nucleon rest frame as:

$$\sin \varphi = \frac{\vec{q} \cdot [(\vec{q} \times \vec{p}_D) \times (\vec{q} \times \vec{k})]}{|\vec{q}| |\vec{q} \times \vec{p}_D| |\vec{q} \times \vec{k}|}, \quad (7)$$

while the final nucleon momentum lies in the xz plane ($\Delta^y = 0$).

We now calculate from T_L and T_T the quantities σ_{00} , σ_{--} and σ_{-0} which enter the differential cross section (6). The longitudinal and transverse cross sections σ_{00} and σ_{--} , at zeroth order in Δ_T , depend on longitudinal $\mathcal{H}_D, \tilde{\mathcal{H}}_D, \tilde{\mathcal{E}}_D$ and transverse form factors through :

$$\sigma_{00} = \frac{C^2}{2Q^2} \left\{ 8(|\mathcal{H}_D^2| + |\tilde{\mathcal{H}}_D^2|)(1 - \xi^2) + |\tilde{\mathcal{E}}_D^2| \frac{1 + \xi^2}{1 - \xi^2} \right\}, \quad (8)$$

$$\sigma_{--} = \frac{4\xi^2 C^2 m_c^2}{Q^4} \left\{ (1 - \xi^2) |\mathcal{H}_T^\phi|^2 + \frac{\xi^2}{1 - \xi^2} |\tilde{\mathcal{E}}_T^\phi|^2 - 2\xi \mathcal{R}e[\mathcal{H}_T^\phi \tilde{\mathcal{E}}_T^{\phi*}] \right\}. \quad (9)$$

The interference cross section σ_{-0} has its first non-vanishing contribution linear in Δ_T/m_N (with $\lambda = \tau^* = 1 + i2$):

$$\begin{aligned} \sigma_{-0} = & \frac{-\xi\sqrt{2}C^2}{m_N} \frac{m_c}{Q^3} \left\{ -i\mathcal{H}_T^{*\phi} \tilde{\mathcal{E}}_D \xi(1+\xi) \epsilon^{pm\Delta\lambda} + \mathcal{H}_T^{*\phi} \Delta^\lambda [-(1+\xi)\mathcal{E}_D] \right. \\ & + \tilde{\mathcal{H}}_T^{*\phi} \Delta^\lambda [2\mathcal{H}_D - \frac{2\xi^2}{1-\xi^2} \mathcal{E}_D] + \mathcal{E}_T^{*\phi} \Delta^\lambda [(1-\xi^2)\mathcal{H}_D - \xi^2 \mathcal{E}_D] \\ & \left. + \tilde{\mathcal{E}}_T^{*\phi} [\Delta^\lambda [(1+\xi)\mathcal{H}_D + \xi\mathcal{E}_D] + i(1+\xi) \epsilon^{pm\Delta\lambda} \tilde{\mathcal{H}}_D] \right\}. \end{aligned} \quad (10)$$

The quantities $\mathcal{R}e(\sigma_{-0})$ and $\mathcal{I}m(\sigma_{-0})$ are directly related to the observables $\langle \cos \varphi \rangle$ and $\langle \sin \varphi \rangle$ through

$$\begin{aligned} \langle \cos \varphi \rangle &= \frac{\int \cos \varphi \, d\varphi \, d^4\sigma}{\int d\varphi \, d^4\sigma} = K_\epsilon \frac{\mathcal{R}e\sigma_{-0}}{\sigma_{00}}, \\ \langle \sin \varphi \rangle &= K_\epsilon \frac{\mathcal{I}m\sigma_{-0}}{\sigma_{00}}, \end{aligned} \quad (11)$$

with $K_\epsilon = \frac{\sqrt{1+\epsilon} + \sqrt{1-\epsilon}}{2\sqrt{\epsilon}}$ and where we consistently neglected the $O(\frac{m_c^2}{Q^2})$ contribution of σ_{--} in the denominator. In the legitimate limit of small α , the dependence on the heavy meson DA effectively disappears in the ratios of the r.h.s. of Eq. (11). The complete formula is quite lengthy but a simple approximate result reads:

$$\begin{aligned} \langle \cos \varphi \rangle &\approx \frac{K \mathcal{R}e[\mathcal{H}_D(2\tilde{\mathcal{H}}_T^\phi + \mathcal{E}_T^\phi + \tilde{\mathcal{E}}_T^\phi)^* - \mathcal{E}_D \mathcal{H}_T^{\phi*}]}{8|\mathcal{H}_D^2| + |\tilde{\mathcal{E}}_D^2|}, \\ \langle \sin \varphi \rangle &\approx \frac{K \mathcal{I}m[\mathcal{H}_D(2\tilde{\mathcal{H}}_T^\phi + \mathcal{E}_T^\phi + \tilde{\mathcal{E}}_T^\phi)^* - \mathcal{E}_D \mathcal{H}_T^{\phi*}]}{8|\mathcal{H}_D^2| + |\tilde{\mathcal{E}}_D^2|}, \\ K &= -\frac{\sqrt{1+\epsilon} + \sqrt{1-\epsilon}}{2\sqrt{\epsilon}} \frac{2\sqrt{2}\xi m_c}{Q} \frac{\Delta_T}{m_N}. \end{aligned} \quad (12)$$

4. Conclusion.

We thus have defined a new way to get access to the transversity chiral-odd generalized parton distributions, the knowledge of which would shed light on the tomography of the quark structure of the nucleons. Planned high energy neutrino facilities [17], which have their scientific program oriented toward the understanding of neutrino oscillations or elusive sterile neutrinos, may thus allow some important progress in the realm of hadronic physics.

This work is partly supported by the Polish Grant NCN No DEC-2011/01/B/ST2/03915 and the French grant ANR PARTONS (Grant No.ANR-12-MONU-0008-01).

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